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Ghost-imaging experiment by measuring reflected photons

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A CCD array is placed facing a chaotic light source and gated by a photon counting detector that simply counts all randomly scattered and reflected photons from an object. A "ghost" image of the object is then observed in the gated CCD. Differing from all published ghost-imaging experiments, this setup captures ghosts from scattered and reflected light of an object, instead of the transmitted ones. This new feature is not only useful for practical applications, but is also important fundamentally. It further explores the nonclassical interference nature of thermal light ghost imaging.

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The first two-photon imaging experiment was demonstrated by Pittman *et al.* in 1995 [1], inspired by the theoretical work of Klyshko [2]. The experiment was immediately named "ghost imaging" due to its surprising nonlocal feature. The important physics demonstrated in that experiment, nevertheless, may not be the "ghost." Indeed, the original purpose of the experiment was to study and to test the two-particle EPR [3] correlation in position and in momentum for an entangled two-photon system [1,4]. The experiments of ghost imaging [1] and ghost interference [5] together stimulated the foundation of quantum imaging (QI) in terms of multiphoton geometrical and physical optics.

Entangled multiphoton systems were later introduced to lithography for sub-diffraction-limited imaging [6]. In 2000, Boto *et al.* proposed a "noon" state and proved that the entangled N-photon system may improve the spatial resolution of an imaging system by a factor of N, despite the Rayleigh diffraction limit. The working principle of quantum lithography was experimentally demonstrated by D'Angelo *et al.* in 2001 [7] by taking advantage of an entangled two-photon state of spontaneous parametric down conversion.

QI has so far demonstrated two peculiar features: (1) reproducing ghost images in a "nonlocal" manner, and (2) enhancing the spatial resolution of imaging beyond the diffraction limit. Both the nonlocal behavior observed in the ghost-imaging experiment and the apparent violation of the uncertainty principle explored in the quantum lithography experiment are due to the two-photon coherent effect of entangled states, which involves the superposition of two-photon amplitudes, a nonclassical entity corresponding to different yet indistinguishable alternative ways of triggering a joint-detection event in the quantum theory of photodetection [8].

In 2004, Gatti *et al.* [9], Wang *et al.* [10], and Zhu *et al.* [11] proposed using thermal radiation to replace the entangled state. A question about ghost imaging is then naturally raised: Is ghost imaging a quantum effect if it can be simulated by "classical" light [12]? Thermal light ghost imaging is based on the second-order spatial correlation of thermal radiation. In fact, two-photon correlation of thermal radiation is not a new observation. Hanbury-Brown and Twiss (HBT) demonstrated the second-order spatial correlation of thermal light in 1956 [13]. Differing from entangled states, the maximum correlation in thermal radiation is 50%, which means 33% visibility of intensity modulation at most. Nevertheless, thermal light is a useful candidate for ghost imaging in practical applications. Recently, a number of experiments successfully demonstrated certain interesting features of ghost imaging by using chaotic light [14–17].

The HBT experiment was successfully interpreted as a statistical correlation of intensity fluctuations. In HBT, the measurement is in the far field (equivalently the Fourier transform plane). The measured two intensities have the same fluctuations while the two photodetectors receive the same mode yielding maximum correlation

$$\langle I_1 I_2\rangle = \bar{I}_1\bar{I}_2 + \langle \Delta I_1 \Delta I_2\rangle. \quad (1)$$

When the two photodetectors receive different modes the intensities have different fluctuations, the measurement yields $\langle \Delta I_1\Delta I_2\rangle=0$ and gives $\langle I_1 I_2\rangle=\bar{I}_1\bar{I}_2$. One type of the HBT experiments explored the partial (50%) spatial correlation of the thermal radiation field, $\langle I_1 I_2\rangle \sim I_0^2\{1+\mathrm{sinc}^2[\pi\Delta\theta(x_1+x_2)/\lambda]\}$, where x_j is the transverse coordinate of the jth photodetector in one-dimension (1D) and $\Delta\theta$ is the angular size of the source. This result has been applied in astronomy for measuring the angular size of stars.

Although Eq. (1) gives a reasonable explanation to the far-field HBT phenomena, it fails to provide an adequate interpretation to a recent lensless ghost-imaging experiment of Scarcelli *et al.* [15]. Differing from HBT in which the measurement is in the far field, Scarcelli's experiment is in the near field. In the near field, for any point on the detection plane, a point photodetector receives a large number of (N) modes in the measurement. The ratio between joint detections triggered by "identical mode" and joint detections triggered by "different modes" is $N/N^2=1/N$. For a large N, the contributions from identical mode are negligible and thus $\langle \Delta I_1 \Delta I_2\rangle=0$, as we know that different modes of chaotic light fluctuate randomly and independently [18]. Therefore, the classical idea of statistical correlation of intensity fluctuations will not work in the multimode case. On the other hand, Scarcelli *et al.* proved a successful alternative interpretation in terms of the quantum theory of two-photon interference.

The two-photon superposition is a new concept that has benefitted from the research of entangled states [4]. How-

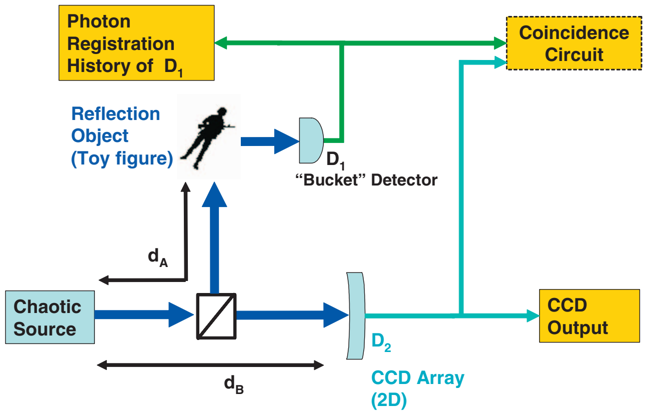

FIG. 1. (Color) ARL ghost image experiment schematic.

ever, the concept of two-photon superposition is not restricted to the entangled states. It is generally true for any radiation, including "classical" thermal light. Unfortunately, this concept has no counterpart in the classical electromagnetic theory of light.

We wish to report an experimental study of near-field thermal light ghost imaging along the same line of Scarcelli *et al.* to provide further experimental evidence and theoretical analysis in supporting the quantum theory of ghost imaging [19]. Differing from Scarcelli's experiment and all other published ghost-imaging demonstrations, this setup captures the ghosts by counting the randomly *scattered* and *reflected* photons from the surface of an object, instead of measuring the transmitted rays. This new feature is not only useful for practical imaging-sensing field applications, but is also important to fundamental concerns. It rejects the classical "projection shadow" model of ghost imaging [20] in a nondeniable way.

Figure 1 is the schematic setup of the experiment. Radiation from a chaotic pseudothermal source [21,22] is divided into two paths by a nonpolarizing beam splitter. In arm A, an object, such as a toy figure, was illuminated by the light source at a distance of $d_A=450$ mm. A "bucket" photodetector, D_1, was used to collect and to count the photons that were reflected from the surface of the object. In arm B a two-dimensional (2D) photon counting CCD array, cooled for single-photon's detection, was placed a distance of $d_B=d_A=450$ mm from the source. The CCD array was facing the light source instead of facing the object. The bucket detector was simulated by using a large area silicon photodiode for collecting the randomly scattered and reflected photons from the object. A triggering pulse from a PC was used to synchronize the measurements at D_1 and the CCD array for two-photon joint detection. The time window was carefully chosen to match the coherent time of the radiation. The light intensity is also carefully chosen for each element of the CCD working at a single-photon level within the period of its response time. The chaotic light was simulated by transmitting a laser beam first through a lens to widen the beam and then through a phase screen made from rotating ground glass. A large transverse sized source gives better spatial resolution of the two-photon image [23].

Figure 2 reports the ghost image of the toy figure. Although the image quality definitely needs to be improved, it is pretty clear what the object is. The poor quality of the image is mainly due to the low photon flux of the reflection.

FIG. 2. Ghost image of a toy figure.

To be sure the new experimental setup is equivalent to that of the historical ghost-imaging experiments, we have also performed a similar measurement. Figure 3 reports a ghost image of an "ARL" name stencil mask. In this measurement, the bucket detector D_1 was placed behind the ARL stencil mask, and collects and counts the transmitted photons that have passed through the ARL letters. The result shows a high fidelity reproduction of the letters "ARL." When the CCD was moved away from $d_B=d_A$, the images were blurred. We are thus sure it is an image and not a "projection shadow."

There is no doubt the toy figure in Fig. 2 is an image by any standard meaning, except the image exists in joint detection only. Mathematically, a perfect ghost image is the result

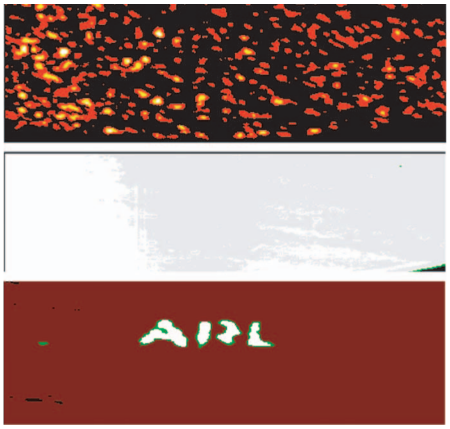

FIG. 3. (Color) Ghost image of "ARL" stencil. Upper: single frame CCD output. The "speckles" indicate typical random photodetection events. This experiment does not have a point-to-point imaging relationship between the source and the CCD. Each speckle of the CCD can be excited by any or by all speckles of the source. Middle: time averaged CCD output of a few hundred frames. Lower: CCD-D_1 joint detection.

of a convolution between the aperture function (amplitude distribution function) of the object $A(\vec{\rho}_o)$ and a δ-function like second-order correlation function $G^{(2)}(\vec{\rho}_o,\vec{\rho}_i)$,

$$F(\vec{\rho}_i) = \int_{obj} d\vec{\rho}_o A(\vec{\rho}_o) G^{(2)}(\vec{\rho}_o,\vec{\rho}_i), \tag{2}$$

where $G^{(2)}(\vec{\rho}_o,\vec{\rho}_i) \simeq \delta(\vec{\rho}_o - \vec{\rho}_i/m)$, $\vec{\rho}_o$ and $\vec{\rho}_i$ are 2D vectors of the transverse coordinate in the object plane and the image plane, respectively, and m is the magnification factor. The δ function characterizes the perfect point-to-point relationship between the object plane and the image plane. If the image comes with a constant background, as in this experiment, the second-order correlation function $G^{(2)}(\vec{\rho}_o,\vec{\rho}_i)$ in Eq. (2) must be composed of two parts

$$G^{(2)}(\vec{\rho}_o,\vec{\rho}_i) = G_0 + \delta(\vec{\rho}_o - \vec{\rho}_i/m), \tag{3}$$

where G_0 is a constant. The value of G_0 determines the visibility of the image. One may immediately connect Eq. (3) with the $G^{(2)}$ function of thermal radiation

$$G^{(2)} = G^{(1)}_{11} G^{(1)}_{22} + |G^{(1)}_{12}|^2, \tag{4}$$

where $G^{(1)}_{11} G^{(1)}_{22} \sim G_0$ is a constant, and $|G^{(1)}_{12}|^2 \sim \delta(\vec{\rho}_1 - \vec{\rho}_2)$ represents a nonlocal position-position correlation. Although the second-order correlation function $G^{(2)}$ is formally written in terms of $G^{(1)}$s, the physics are completely different. As we know, $G^{(1)}_{12}$ is usually measured by one photodetector representing the first-order coherence of the field, i.e., the ability of observing first-order interference. Here, in Eq. (4), $G^{(1)}_{12}$ is measured by two independent photodetectors at distant space-time points and represents a nonlocal EPR correlation.

Differing from the phenomenological classical theory of intensity-intensity correlation, the quantum theory of joint photodetection (Glauber's theory) [8] dips into the physical origin of the phenomenon. The theory gives the probability of a specified joint photodetection event

$$G^{(2)} = \mathrm{Tr}[\hat{\rho} E^{(-)}(\vec{\rho}_1) E^{(-)}(\vec{\rho}_2) E^{(+)}(\vec{\rho}_2) E^{(+)}(\vec{\rho}_1)], \tag{5}$$

and leaves room for us to identify the superposed probability amplitudes. In Eq. (5), $E^{(-)}$ and $E^{(+)}$ are the negative and positive-frequency field operators at space-time coordinates of the photodetection event and $\hat{\rho}$ represents the density operator describing the radiation. In Eq. (5), we have simplified the calculation to 2D.

In the *photon counting* regime, it is reasonable to model the thermal light in terms of *single-photon states* for joint dection (see the Appendix),

$$\hat{\rho} \simeq |0\rangle\langle 0| + |\epsilon|^4 \sum_{\vec{\kappa}} \sum_{\vec{\kappa}'} \hat{a}^\dagger(\vec{\kappa}) \hat{a}^\dagger(\vec{\kappa}') |0\rangle\langle 0| \hat{a}(\vec{\kappa}') \hat{a}(\vec{\kappa}), \tag{6}$$

where $|\epsilon| \ll 1$. Basically, we model the state of thermal radiation, which results in a joint-detection event, as a statistical mixture of two photons with equal probability of having any transverse momentum $\vec{\kappa}$ and $\vec{\kappa}'$.

The second-order transverse spatial correlation function is thus

$$G^{(2)}(\vec{\rho}_1,\vec{\rho}_2) = \sum_{\vec{\kappa},\vec{\kappa}'} |\langle 0| E^{(+)}_2(\vec{\rho}_2) E^{(+)}_1(\vec{\rho}_1) |1_{\vec{\kappa}} 1_{\vec{\kappa}'}\rangle|^2. \tag{7}$$

The electric field operator, in terms of the transverse mode and coordinates, can be written as follows:

$$E^{(+)}_j(\vec{\rho}_j) \propto \sum_{\vec{\kappa}} g_j(\vec{\kappa};\vec{\rho}_j) \hat{a}(\vec{\kappa}), \tag{8}$$

where $\hat{a}(\vec{\kappa})$ is the annihilation operator for the mode corresponding to $\vec{\kappa}$ and $g_j(\vec{\rho}_j;\vec{\kappa})$ is the Green's function associated with the propagation of the field from the source to the jth detector [24]. Substituting the field operators into Eq. (7), we obtain

$$G^{(2)}(\vec{\rho}_1,\vec{\rho}_2) = \sum_{\vec{\kappa},\vec{\kappa}'} |g_2(\vec{\kappa};\vec{\rho}_2) g_1(\vec{\kappa}';\vec{\rho}_1) + g_2(\vec{\kappa}';\vec{\rho}_2) g_1(\vec{\kappa};\vec{\rho}_1)|^2. \tag{9}$$

Equation (9) indicates a two-photon superposition. The superposition happens between two different yet indistinguishable Feynman alternatives that lead to a joint photodetection: (1) photon $\vec{\kappa}$ and photon $\vec{\kappa}'$ are annihilated at $\vec{\rho}_2$ and $\vec{\rho}_1$, respectively, and (2) photon $\vec{\kappa}'$ and photon $\vec{\kappa}$ are annihilated at $\vec{\rho}_2$ and $\vec{\rho}_1$, respectively. The interference phenomenon is not, as in classical optics, due to the superposition of electromagnetic fields at a local point of space time. It is due to the superposition of $g_2(\vec{\kappa};\vec{\rho}_2) g_1(\vec{\kappa}';\vec{\rho}_1)$, and $g_2(\vec{\kappa}';\vec{\rho}_2) g_1(\vec{\kappa};\vec{\rho}_1)$, the so-called two-photon amplitudes.

Completing the normal square of Eq. (9), it is easy to find that the sum of the normal square terms corresponding to the constant of G_0 in Eq. (3): $\Sigma_{\vec{\kappa}} |g_1(\vec{\kappa};\vec{\rho}_1)|^2 \Sigma_{\vec{\kappa}'} |g_2(\vec{\kappa}';\vec{\rho}_2)|^2 = G^{(1)}_{11} G^{(1)}_{22}$, and the cross term $|\Sigma_{\vec{\kappa}} g^*_1(\vec{\kappa};\vec{\rho}_1) g_2(\vec{\kappa};\vec{\rho}_2)|^2 = |G^{(1)}_{12}(\vec{\rho}_1,\vec{\rho}_2)|^2$ gives the δ function of position-position correlation

$$\left| \int d\vec{\kappa} g^*_1(\vec{\kappa};\vec{\rho}_1) g_2(\vec{\kappa};\vec{\rho}_2) \right|^2 \simeq |\delta(\vec{\rho}_o + \vec{\rho}_i)|^2, \tag{10}$$

where

$$g_1(\vec{\kappa};\vec{\rho}_o) \propto \Psi\left(\vec{\kappa}, -\frac{c}{\omega} d_A\right) e^{i\vec{\kappa}\cdot\vec{\rho}_o},$$

$$g_2(\vec{\kappa};\vec{\rho}_i) \propto \Psi\left(\vec{\kappa}, -\frac{c}{\omega} d_B\right) e^{i\vec{\kappa}\cdot\vec{\rho}_i}, \tag{11}$$

are the Green's functions propagated from the radiation source to the transverse planes of d_A and $d_B = d_A$. In Eq. (11), $\Psi(\omega d/c)$ is a phase factor representing the optical transfer function of the linear system under the Fresnel near-field paraxial approximation, ω is the frequency of the radiation field, and c is the speed of light.

Substituting this δ function together with the constant G_0 into Eq. (2), an equal sized lensless image of $A(\vec{\rho}_o)$ is observed in the joint detection between the CCD array and the photon counting detector D_1. The visibility of the image is determined by the value of G_0.

The experiment is thus successfully interpreted as the result of two-photon interference. The two-photon interference results in a point-point correlation between the object plane and the image plane and yields a ghost image of the object by means of joint photodetection.

In summary, we have demonstrated a new type of ghost-imaging experiment. This interesting experiment is useful for practical imaging-sensing field applications, and for the fundamental understanding of the nonlocal ghost imaging phenomenon as well as the quantum mechanical concept of multiphoton superposition.

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APPENDIX: QUANTUM STATE OF THERMAL LIGHT

We assume a large number of atoms that are ready for two-level atomic transition. At most times, the atoms are in their ground state. There is, however, a small chance for each atom to be excited to a higher energy level and later release a photon during an atomic transition from the higher energy level E_2 ($\Delta E_2 \neq 0$) back to the ground state E_1. It is reasonable to assume that each atomic transition excites the field into the following state:

$$|\Psi\rangle \simeq |0\rangle + \epsilon \sum_{\mathbf{k},s} f(\mathbf{k},s)\hat{a}^{\dagger}_{\mathbf{k},s}|0\rangle,$$

where $|\epsilon| \ll 1$ is the probability amplitude for the atomic transition. Within the atomic transition, $f(\mathbf{k},s) = \langle \Psi_{\mathbf{k},s}|\Psi\rangle$ is the probability amplitude for the radiation field to be in the single-photon state of wave number $\mathbf{k}$ and polarization $s: |\Psi_{\mathbf{k},s}\rangle = |1_{\mathbf{k},s}\rangle = \hat{a}^{\dagger}_{\mathbf{k},s}|0\rangle$.

For this simplified two-level system, the density matrix that characterizes the state of the radiation field excited by a large number of possible atomic transitions is thus

$$\begin{aligned}\hat{\rho} &= \prod_{t_{0j}} \left\{ |0\rangle + \epsilon \sum_{\mathbf{k},s} f(\mathbf{k},s) e^{-i\omega t_{0j}} \hat{a}^{\dagger}_{\mathbf{k},s}|0\rangle \right\} \\ &\quad \times \prod_{t_{0k}} \left\{ \langle 0| + \epsilon^* \sum_{\mathbf{k}',s'} f(\mathbf{k}',s') e^{i\omega' t_{0k}} \langle 0|\hat{a}_{\mathbf{k}',s'} \right\} \\ &\simeq \left\{ |0\rangle + \epsilon \left[\sum_{t_{0j}} \sum_{\mathbf{k},s} f(\mathbf{k},s) e^{-i\omega t_{0j}} \hat{a}^{\dagger}_{\mathbf{k},s}|0\rangle \right] + \epsilon^2[\cdots] \right\} \\ &\quad \times \left\{ \langle 0| + \epsilon^* \left[\sum_{t_{0k}} \sum_{\mathbf{k}',s'} f(\mathbf{k}',s') e^{i\omega' t_{0k}} \langle 0|\hat{a}_{\mathbf{k}',s'} \right] \right. \\ &\quad \left. + \epsilon^{*2}[\cdots] \right\},\end{aligned}$$

where $e^{-i\omega t_{0j}}$ is a random phase factor associated with the state $|\Psi_j\rangle$ of the jth atomic transition. Summing over t_{0j} and t_{0k} by taking all possible values, we find the approximation to the fourth order of $|\epsilon|$,

$$\begin{aligned}\hat{\rho} &\simeq |0\rangle\langle 0| + |\epsilon|^2 \sum_{\mathbf{k,s}} |f(\mathbf{k,s})|^2 |1_{\mathbf{k,s}}\rangle\langle 1_{\mathbf{k,s}}| \\ &\quad + |\epsilon|^4 \sum_{\mathbf{k,s}} \sum_{\mathbf{k}',s'} |f(\mathbf{k,s})|^2 |f(\mathbf{k}',s')|^2 |1_{\mathbf{k,s}} 1_{\mathbf{k}',s'}\rangle\langle 1_{\mathbf{k},s} 1_{\mathbf{k}',s'}|.\end{aligned}$$

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